RESONANT NEUTRINO SPIN–FLAVOR PRECESSION AND SUPERNOVA SHOCK REVIVAL

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Abstract

A new mechanism of supernova shock revival is proposed, which involves resonant spin–flavor precession of neutrinos with a transition magnetic moment in the magnetic field of the supernova. The mechanism can be operative in supernovae for transition magnetic moments as small as $10^{-14}\mu_B$ provided the neutrino mass squared difference is in the range $\Delta m^2 \sim (3 \text{ eV})^2 - (600 \text{ eV})^2$. It is shown that this mechanism can increase the neutrino–induced shock reheating energy by about 60%.

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1 Introduction

One of the most important problems in the theory of supernova explosions is to understand the physical mechanisms which eventually expel the outer mantle delivering the right amount of energy. Although the main ideas involved in the theory have received major confirmation after the detection of the neutrinos from SN1987A, the problem of accelerating the outward going shock wave which forms after core collapse is still unresolved. For many years, when all computer calculations were essentially done in one dimension, the majority of computations were unsuccessful since the shock would travel for about $300 - 500$ km and then stall, after losing energy by dissociating heavy nuclei in the envelope into nucleons. Many proposals have been made to solve this problem, starting from including general relativistic corrections, different equations of state, better neutrino transport description and new neutrino physics. Recent numerical calculations [1, 2, 3] have shown that as one moves to more than one dimension one can easily get convective instabilities driven by neutrino heating which are very effective and fast in reheating the material behind the shock. The idea that convection would help the explosion is not new. It has been in the literature for some time, see, e.g., [4] (see [2] for an extensive list of references). The idea that multi-dimensional calculations might help is also not new. Several attempts have been made to model supernova explosion in more than one dimension; probably they failed because they did not contain the right combinations of other aspects of the physics involved. Although these most recent calculations make an important contribution to the subject we still believe that much work should be done to understand fully the mechanisms involved. The proposed convective instability relies heavily on the details of neutrino interactions with matter which control the energy transport process. The details of this process are still controversial, especially when the matter is at high density and is not spherically symmetrical. The question whether convective instabilities can revive the shock and lead to a successful supernova explosion is therefore still far from settled, and any new mechanism which could contribute towards the shock energy would be very welcome.

Recently there has been discussion of whether massive neutrinos, which seem to be necessary to reconcile the solar neutrino experiments with the standard model of the sun, might also help in accelerating the shock. Matter–enhanced neutrino oscillations (MSW effect, [5, 6, 7]) in supernovae might play an important role in reviving the shock after core collapse by increasing the amount of energy that neutrinos deposit behind the shock [8, 9, 10]. The idea is based on the fact that in the region between the neutrinosphere and the position of the stalled shock the matter density is such that flavor transformation of ν_μ or ν_τ into ν_e is resonant for masses of the heavier neutrinos in the range $10 - 100$ eV; this transformation can be efficient even if the vacuum neutrino mixing angle is quite small, $\theta \gtrsim 10^{-4}$. Since the average energy of ν_μ's and ν_τ's at the neutrinosphere is about 20 MeV whereas that of ν_e's is about 10 MeV, the electron neutrinos emerging as a result of the $\nu_\mu(\nu_\tau) \to \nu_e$ transformation would have twice as high an energy as the originally emitted ones, and this extra energy would be available for heating the matter behind the shock. Electron neutrinos interact with matter with a larger cross section than muon or tauon neutrinos since they have charged–current interactions with matter in addition to the neutral–current ones. Therefore the ν_e's produced in the $\nu_\mu(\nu_\tau) \to \nu_e$ transformation will more efficiently deposit energy behind the shock. Fuller et al. [10] have shown that the net effect is a $\sim 60\%$ increase in the supernova explosion energy.

In this paper, we show that a similar result can be obtained in the framework of the resonant spin–flavor precession mechanism. In this case one has to assume that the neutrino has a nonzero transition magnetic moment μ by which it interacts with a magnetic field. Since we know that after a supernova explosion a pulsar is, in many cases, left over, the important role played by the magnetic field during and after core collapse is not in question. The strong magnetic field and the high density make the environment between the neutrinosphere and the position of the stalled shock suitable for spin–flavor conversion due to transition magnetic moments of neutrinos.

Very much as in the case of the MSW effect, spin–flavor precession due to a transition

magnetic moment of neutrinos, in which neutrino helicity and flavor are rotated simultaneously [11], can be resonantly enhanced in matter [12, 13]. This effect can explain the observed deficit of solar neutrinos with respect to the predictions of the standard solar model [12, 13, 14, 15, 16, 17]. That requires the neutrino transition magnetic moment to be of the order of $\mu \approx 10^{-11}\mu_B$ ($\mu_B = e/2m_e$ is the electron Bohr magneton) provided the strength of the magnetic field near the bottom of the convective zone of the sun is of the order of a few tens of kG. The indicated value of the transition magnetic moment is to be compared with recent astrophysical upper bounds derived from the limits on the energy loss rates of white dwarfs ($10^{-11}\mu_B$, ref. [18]) and helium stars ($3 \times 10^{-12}\mu_B$, ref. [19], and $10^{-12}\mu_B$, ref. [20]). The latter two values imply that an order of magnitude stronger magnetic field might be necessary to account for the solar neutrino problem in the framework of the neutrino magnetic moment scenario.

In the present paper we show that the spin–flavor precession of neutrinos may play an important role in supernova dynamics even if neutrino transition magnetic moments are far below the present astrophysical upper limits. In particular, it can be resonantly enhanced in the region between the neutrinosphere and the position of the shock for typical values of $\mu \approx 10^{-14}\mu_B$, magnetic field strengths of $B \approx 10^{12} - 10^{15}$ G and neutrino masses which lie in the range $\sim (3 - 600)$ eV. This strength of the magnetic field is natural in the context of supernovae if the explosion does give rise to a pulsar, the value of μ is consistent with the prediction of the decaying neutrino hypothesis [21], and the range of neutrino masses is the one which is relevant for cosmology and the decaying neutrino theory.

The idea that neutrino magnetic moments can play an important role in supernova dynamics was first put forward by Dar [22] in the context of the usual Dirac neutrino magnetic moments and transitions of active left-handed neutrinos into sterile right-handed ones. Our mechanism, based on the spin–flavor precession of neutrinos due to their Majorana-like transition magnetic moments, is different from Dar's. Resonant spin–flavor precession (RSFP) of neutrinos in type II supernovae has been studied earlier [23, 24]; however the main goal of

those papers was to explore possible consequences for the neutrino signal from supernovae, and no implications of this effect for supernova dynamics were discussed.

The plan of the paper is as follows. We start by reviewing the main features of RSFP in Sec. 2. In Sec. 3 we discuss the RSFP in supernovae, and in Sec. 4 consider the implications of this effect for supernova shock reheating. Sec. 5 contains our conclusions.

2   Basic features of RSFP

We confine our analysis to the simplest case of Majorana neutrinos, for which the diagonal magnetic moments are zero. We also assume that there are only two neutrino flavors and disregard the usual flavor mixing, i.e. we consider transitions in a two-neutrino system ν_e and ν_a (ν_a can be either ν_τ or ν_μ) with masses m_{ν_e} and m_{ν_a} due to the transition magnetic moment $\mu_{ea} \equiv \mu$. More specifically, of main interest to us are the transitions between the right–handed electron antineutrinos $\bar{\nu}_{eR}$ and left–handed muon or tauon neutrinos $\nu_{\mu L}$ or $\nu_{\tau L}$ (ν_{aL}). As we shall see, the transitions between the corresponding antiparticles, namely $\nu_{eL} \leftrightarrow \bar{\nu}_{\mu R}(\bar{\nu}_{\tau R})$, which may be relevant for the solution of the solar neutrino problem, are non-resonant in the supernova environment behind the shock provided $m_{\nu_a} > m_{\nu_e}$. We shall comment on the case $m_{\nu_a} < m_{\nu_e}$ in Sec. 4.

In a medium mainly composed of electrons, neutrons and protons the evolution of the $\bar{\nu}_{eR}$ and ν_{aL} states in a transverse magnetic field $B_\perp$ is described by the Schrödinger–like equation [12, 13]

$$i\frac{d}{dt}\begin{pmatrix} \bar{\nu}_{eR} \\ \nu_{aL} \end{pmatrix} = \begin{pmatrix} -A & B \\ B & A \end{pmatrix} \begin{pmatrix} \bar{\nu}_{eR} \\ \nu_{aL} \end{pmatrix} \tag{1}$$

where $A = \Delta m^2/4E + V/2$, $B = \mu B_\perp$, $\Delta m^2 \equiv m_{\nu_a}^2 - m_{\nu_e}^2$ and V is the difference of the effective potentials that neutrinos ν_{aL} and $\bar{\nu}_{eR}$ experience in matter. The effective potential V can be written as a sum of two contributions, $V = V^{(1)} + V^{(2)}$. Numerically the most important term $V^{(1)}$ is due to the interaction of neutrinos with matter: $V^{(1)} =$

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$\sqrt{2}G_F(N_e - N_n)$ where G_F is the Fermi constant and N_e and N_n are the electron and neutron number densities respectively. The second term is due to neutrino–neutrino forward scattering, which in the case of scattering of an electron neutrino by the neutrino sea of all types takes the form [25, 26, 27]

$$V^{(2)}(\nu_e) = \sqrt{2}G_F(2N_{\nu_e}^{eff} + N_{\nu_\mu}^{eff} + N_{\nu_\tau}^{eff}), \qquad (2)$$

where $N_{\nu_e}^{eff}$ is the difference between the effective number densities of electron neutrinos and antineutrinos, and similarly for ν_μ and ν_τ. The word "effective" is used here because the neutrino densities are modified by the averaged factor $\langle 1 - v_m \cos \alpha \rangle$ which takes into account the dependence of the corresponding contributions to the effective potential on the angle α between the momenta of interacting neutrinos. This effect is negligible for neutrino scattering on non-relativistic or randomly-moving particles, but leads to a strong suppression of the $V^{(2)}$ potential for neutrino scattering on relativistic neutrinos ($v_m = 1$) moving nearly in the same direction [27, 10, 28]. Expressions similar to eq. (2) hold also for the effective potentials of μ and τ neutrinos. Antineutrinos will experience the opposite sign potential, $V^{(2)}(\bar{\nu}_l) = -V^{(2)}(\nu_l)$, $l = e, \mu, \tau$. In the case of $\nu_{\mu(\tau)L} \leftrightarrow \bar{\nu}_{eR}$ transitions, the effective potential V reduces to

$$V = \sqrt{2}G_F[N_e - N_n + 3N_{\nu_e}^{eff} + 5N_{\nu_\mu}^{eff}], \qquad (3)$$

where we have taken into account that in the supernova environment $N_{\nu_\mu}^{eff} = N_{\nu_\tau}^{eff}$. However, as we shall see, the neutrino-neutrino forward scattering contributions to V are very small in the shock reheating epoch ($t \simeq 0.15$ s after the shock bounce) and can be safely neglected. Typical average neutrino energies at this epoch are [10] $\langle E_{\nu_e} \rangle \approx 9$ MeV, $\langle E_{\bar{\nu}_e} \rangle \approx 12$ MeV and $\langle E_{\nu_\mu} \rangle \approx \langle E_{\nu_\tau} \rangle \approx 20$ MeV. The neutrinosphere is at $R_\nu \approx 50$ km and the shock position is at about 400 km. The neutrino luminosities at the shock reheating epoch are $L_{\nu_e} \approx L_{\bar{\nu}_e} \approx L_{\nu_{\tau(\mu)}} \approx L_{\bar{\nu}_{\tau(\mu)}} \approx 5 \times 10^{52}$ erg s^{-1}. By inserting these values into the expression for the effective number density (see eq. (5) in [28]) one gets

$$N_{\nu_e}^{eff} = 1.44 \times 10^{34} \text{ cm}^{-3} \left(\frac{10 \text{ km}}{r} \right)^4, \quad N_{\nu_\mu}^{eff} = N_{\nu_\tau}^{eff} = 0, \qquad (4)$$

and therefore the effective potential due to neutrino–neutrino scattering in the case of $\nu_{\mu(\tau)L} \leftrightarrow \bar{\nu}_{eR}$ transitions is

$$V^{(2)} \approx -5.45 \times 10^{-3} \left(\frac{10 \text{ km}}{r} \right)^4 \text{ eV}. \tag{5}$$

This term is numerically very small compared to the main $V^{(1)}$ term in the region of interest to us, namely $50 \text{ km} \lesssim r \lesssim 400 \text{ km}$.

From eq. (1) one can find the resonance condition by equating the diagonal terms of the effective Hamiltonian of the neutrino system:

$$\sqrt{2} G_F \frac{\rho}{m_N} (1 - 2Y_e) = \frac{\Delta m^2}{2E}, \tag{6}$$

or

$$7.54 \times 10^{-14} \text{ eV } \rho(\frac{\text{g}}{\text{cc}}) \left(1 - 2Y_e \right) = \frac{\Delta m^2}{2E}, \tag{7}$$

where ρ is the matter mass density, m_N is the mass of the nucleon, and Y_e is the number of electrons per baryon. For a given Δm^2 and energy E eq. (6) gives the values of the density at which the transition is resonant. Since in the region between the neutrinosphere and the shock position Y_e is always less than $1/2$ at the shock reheating epoch [10, 28], from eqs. (6) or (7) it follows that for $\Delta m^2 > 0$ only the transitions $\nu_{\mu(\tau)L} \leftrightarrow \bar{\nu}_{eR}$ will be resonant, whereas the transitions between the corresponding antiparticles, for which the signs of the l.h.s. of eqs. (6) and (7) must be reversed, are non-resonant.

The neutrino eigenstates in matter and a magnetic field are linear combinations of $\bar{\nu}_{eR}$ and $\nu_{\mu(\tau)L}$ with the mixing angle defined through

$$\tan 2\theta = \frac{2\mu B_\perp}{-\sqrt{2} G_F (N_n - N_e) + \frac{\Delta m^2}{2E}}. \tag{8}$$

The efficiency of the $\nu_{\mu(\tau)L} \to \bar{\nu}_{eR}$ transition is determined by the degree of the adiabaticity which depends on both the neutrino energy and the magnetic field strength at the resonance:

$$\gamma \equiv \pi \frac{\Delta r}{l_r} = 8 \frac{E}{\Delta m^2} (\mu B_{\perp r})^2 (L_\rho)_r. \tag{9}$$

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Here Δr is the resonance width, $l_r = \pi/\mu B_{\perp r}$ is the precession length at the resonance and $L_\rho \equiv |\frac{1}{\rho(1-2Y_e)}\frac{d\rho(1-2Y_e)}{dr}|^{-1}$ is the characteristic length over which the effective matter density $(1-2Y_e)\rho$ varies significantly in the supernova, $(L_\rho)_r$ being its value at the resonance. For the RSFP to be efficient, γ should be $\gtrsim 1$.

The probability of the $\nu_{\mu(\tau)L} \leftrightarrow \bar{\nu}_{eR}$ transition can be written in the following general form [29]:

$$P(\nu_{\mu(\tau)} \to \bar{\nu}_e) = \frac{1}{2} - \frac{1}{2}\cos 2\theta_i \cos 2\theta_f(1-2P') - \sqrt{P'(1-P')}\cos 2\theta_i \sin 2\theta_f$$
$$\times \cos(\Phi_{12} + \Phi_{22}) + \sqrt{P'(1-P')}\sin 2\theta_i \cos 2\theta_f \ \cos(\Phi_{12} - \Phi_{22})$$
$$+\frac{1}{2}P'\sin 2\theta_i \sin 2\theta_f \ (\cos 2\Phi_{12} + \cos 2\Phi_{22}) - \frac{1}{2}\sin 2\theta_i \sin 2\theta_f \ \cos 2\Phi_{22}. \tag{10}$$

Here θ_i and θ_f are the values of the neutrino mixing angle in matter at the initial and final points of the neutrino path in the magnetic field, which we will assume to be located at the surface of the neutrinosphere and far beyond the resonance at a density much smaller than the resonance density, Φ_{12} and Φ_{22} are two phases which are responsible for the possible oscillatory dependence of the probability (10) on $E/\Delta m^2$, the magnetic field strength and matter density profiles, and P' is the so-called "jump" probability , i.e., the probability of transition of one of the matter eigenstate neutrinos into another in the course of neutrino propagation in the magnetic field of the supernova. The value of the jump probability is therefore a measure of violation of the adiabaticity of neutrino propagation. As we shall argue in the next Section, for the values of $E/\Delta m^2$ and the magnetic field strengths of interest to us, $\theta_i \approx \pi/2$ and $\theta_f \approx 0$. Under these conditions the general expression for the transition probability (10) simplifies considerably and reduces to

$$P(\nu_{\mu(\tau)} \to \bar{\nu}_e) \cong \frac{1}{2} - \frac{1}{2}\cos 2\theta_i \cos 2\theta_f(1-2P'). \tag{11}$$

For $P' \cong 0$ (adiabtic transitions) the transition probability can be very close to 1. For the purposes of our further analysis the jump probability can be approximated by the Landau-Zener probability:

$$P' \approx P_{LZ} = \exp(-\frac{\pi}{2}\gamma). \tag{12}$$

3 RSFP in supernovae

We know very little about the magnetic field in the supernova environment. It is certain that at least in those cases in which the explosion leaves a pulsar, the magnetic fields may be as strong as $10^{12} - 10^{14}$ G. However, also in these cases nothing is known about the spatial distribution of the field. We shall assume rather arbitrarily that the radial dependence of the magnetic field strength above the neutrinosphere has the power-law behavior

$$B_\perp(r) = B_0 \left(\frac{r_0}{r}\right)^k , \quad r \geq r_0 , \tag{13}$$

where r_0 is the radius of the neutrinosphere, B_0 is the strength of the field at the neutrinosphere, and $k = 2$ or 3.

For a given $\Delta m^2/E$, the resonance condition (7) determines the value of the matter density (and, given the density profile, of the radial coordinate) at which the transition is most efficient. In order for the resonance to occur between the neutrinosphere which is located at about 50 km and the position of the stalled shock ~ 400 km, $\Delta m^2/E$ should be in the range 5.5×10^{-7} to 2×10^{-2} eV. For typical neutrino energies $\langle E_{\nu_\mu} \rangle \approx \langle E_{\nu_\tau} \rangle \approx 20$ MeV this would correspond to values of Δm^2 in the range 11 eV^2 to 4×10^5 eV^2. If we assume the hierarchical pattern of neutrino masses, this would mean that the mass of the heavier neutrino should lie in the range ~ 3 eV to 600 eV. If the shock stalls at larger distances from the supernova core, even smaller neutrino masses would do.

For the above values of $\Delta m^2/E$ one can readily calculate the magnitudes of the product $\mu B_\perp$ which are necessary in order to have an adiabatic or weakly nonadiabatic transition (i.e. $\gamma \geq 1$). The values of L_ρ at different resonance positions can be read off from Fig. 1. One then gets from eq. (9)

$$\mu B_{\perp r} \geq 3.5 \times 10^{-7} \text{ eV} \tag{14}$$

if the resonance takes place at 50 km, or

$$\mu B_{\perp r} \geq 1.1 \times 10^{-9} \text{ eV} \tag{15}$$

when the resonance takes place at 350 km. Assuming the transition magnetic moment $\mu = 10^{-14}$ $(10^{-12})\mu_B$, this gives the following lower bounds on the magnetic field strength at the resonance:

$$B_{\perp r} \geq 6.0 \times 10^{15} \ (6.0 \times 10^{13}) \ \text{G} \qquad (r_{res} \simeq 50 \ \text{km});$$

$$B_{\perp r} \geq 1.9 \times 10^{13} \ (1.9 \times 10^{11}) \ \text{G} \qquad (r_{res} \simeq 350 \ \text{km}). \qquad (16)$$

Alternatively, if one assumes the magnetic field strength at the neutrinosphere $B_0 = 5 \times 10^{14}$ G and $k = 2$, eqs. (14) and (15) transform into the following lower limit for the transition magnetic moment μ:

$$\mu \geq 10^{-14} \ \text{to} \ 10^{-13}\mu_B, \qquad (17)$$

which is one to two orders of magnitude below the current astrophysical upper limits.

It follows from eqs. (8) and (13) that if the resonance occurs not too close to the neutrinosphere ($r_{res} \gtrsim 60$ km) and $\mu B_0 \ll 10^{-2}$ eV (which, e.g., for $\mu = 10^{-12}\mu_B$ is fulfilled provided that $B_0 \leq 10^{17}$ G), $\theta_i \approx \pi/2$. At the same time, we find that for

$$\mu B_0 < 2.9 \times 10^{-6} \ \text{eV}, \qquad (18)$$

the mixing angle close to the position of the stalled shock becomes very small, i.e. $\theta_f \approx 0$. The bound (18) was obtained assuming the resonance takes place at $r_{res} = 350$ km; for smaller r_{res} it gets relaxed. Comparing the upper bound (18) with the lower bound $\mu B_0 > 5.4 \times 10^{-8}$ eV which follows from eq. (15) for $k = 2$, we see that they do not contradict each other. If the condition (18) is not satisfied, the efficiency of the RSFP transition decreases, but for the adiabatic transitions ($\gamma \gg 1$) the transition probability never goes below $1/2$. For $\theta_i \approx \pi/2$, $\theta_f \approx 0$ the transition probability is essentially determined by the degree of the adiabaticity of the transition.

In Fig. 2 we plot the transition probability P vs $E/\Delta m^2$ for two different magnetic field configurations of eq. (13), with $k = 2$ and $k = 3$. The parameter μB_0 was chosen in each case in such a way as to have the transition from the non-adiabatic to the adiabatic regime

for the range of values of $E/\Delta m^2$ which corresponds to the resonance position between the neutrinosphere and the stalled shock. In what follows we shall assume that the adiabaticity condition (9) and the condition (18) are satisfied, i.e. that the transition $\nu_{\mu(\tau)L} \to \bar{\nu}_{eR}$ is nearly complete.

4   Shock reheating

After the bounce the energy of the shock is dissipated by the dissociation of nuclei, and the shock, after traveling up for some distance (which, following Fuller et al. [10], we assume to be about 400 km), stalls. Meanwhile neutrinos diffuse out from the region where they were trapped; interacting with the matter behind the shock they deliver their energy in the neighborhood. The problem of shock reheating by neutrino interactions was discussed in detail by Bethe and Wilson [30] (hereafter BW85). They considered the neutrino capture processes

$$\nu_e + n \to p + e^-, \quad \bar{\nu}_e + p \to n + e^+ \tag{19}$$

which the electron neutrinos and antineutrinos may undergo after diffusion from the supernova core; no neutrino flavor or spin–flavor transformation was taken into account. However the energy of the electron neutrinos is not sufficient to re-accelerate the shock. The possibility that neutrino flavor conversion due to the MSW effect may result in electron neutrinos having higher energies than that of the originally produced ones, thus increasing the amount of energy they deposit, was pointed out in [8, 9] and then considered in detail in [10]. The authors of the latter paper came to the conclusion that the MSW effect in supernovae can give a net gain in the deposited energy of about 60%. In this Section we estimate the gain in the neutrino–induced shock reheating energy due to the spin–flavor conversion of neutrinos.

As discussed above, as a result of RSFP the μ- and τ- type neutrinos which are more energetic than electron-type neutrinos will be converted into electron antineutrinos. These may then interact with the matter behind the shock more efficiently and deliver more energy

than the originally produced $\bar{\nu}_e$'s. Here we estimate the energy gain due to the RSFP of supernova neutrinos following the consideration performed in [10] for the MSW effect.

The neutrino absorption coefficients due to the reactions (19) can be written as

$$K_i(E_\nu) = N_A Y_i \langle \sigma(E_\nu) \rangle, \tag{20}$$

where N_A is Avogadro's number, Y_i ($i = p$, n) is the appropriate nucleon number per baryon, and $\langle \sigma(E_\nu) \rangle$ is the reaction cross section averaged over the neutrino spectrum. Notice that the cross sections of reactions (19) are essentially quadratic in neutrino energies and therefore $\langle \sigma(E_\nu) \rangle$ for the first and the second reactions is quadratic in the temperatures of electron neutrinos and antineutrinos, respectively. In what follows we will assume that the RSFP is adiabatic for all the neutrino energies of interest ($\gamma \gg 1$), and so the conversion does not distort the neutrino spectra and just leads to their interchange. Assuming that the neutrino spectra are quasi-blackbody, the absorption coefficients can be approximated as [30]

$$K_i(T_\nu) \approx (3.8 \times 10^{-19}\ \mathrm{cm^2\,g^{-1}})\, Y_i T_\nu^2. \tag{21}$$

This takes into account the fact that there are actually fewer neutrinos with high energies than in the genuine blackbody spectrum.

The rate at which the specific energy is deposited behind the shock is then [30]

$$\dot{E}_{BW85} \approx (4\pi R_m)^{-2}[K_n(T_{\nu_e})L_{\nu_e} + K_p(T_{\bar{\nu}_e})L_{\bar{\nu}_e}] - 4\pi j(T_m), \tag{22}$$

where L_{ν_e} and $L_{\bar{\nu}_e}$ are the total ν_e and $\bar{\nu}_e$ luminosities, R_m and T_m the radius and the temperature of the element of matter behind the shock, and $j(T_m)$ is the neutrino emissivity per steradian of the element of matter considered. The last (negative) term in (22) is negligible provided $T_m \ll T_{\nu_e}$. Following [30] and [10], in order to get a rough estimate of the effect we will first assume that this is the case; we will come back to the discussion of the emissivity term later.

The specific energy rate (22) derived in [30] does not take into account possible neutrino conversions. It is, however, straightforward to estimate the gain in energy deposited by

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neutrinos behind the shock due to the RSFP transition $\nu_{\mu(\tau)L} \to \bar{\nu}_{eR}$. The ratio of the specific heating rates with and without RSFP transitions is

$$R_{RSFP} \equiv \frac{\dot{E}_{RSFP}}{\dot{E}_{BW85}} \approx \frac{Y_n + Y_p \left(\frac{T_{\nu_\mu}}{T_{\nu_e}}\right)^2}{Y_n + Y_p \left(\frac{T_{\bar{\nu}_e}}{T_{\nu_e}}\right)^2}, \tag{23}$$

where we have taken into account that the total luminosities of all the neutrino species are approximately equal and that the cross sections of the two reactions in eq. (19) for a given neutrino (antineutrino) energy E_ν are practically the same provided that $E_\nu \gg (m_n - m_p) \approx$ 1.3 MeV.

At the neutrino reheating epoch, $\langle E_{\nu_e} \rangle \approx 9$ MeV, $\langle E_{\bar{\nu}_e} \rangle \approx 12$ MeV and $\langle E_{\nu_\mu} \rangle \approx \langle E_{\nu_\tau} \rangle \approx$ 20 MeV [31]. Substituting these values into eq. (23) and assuming that the reheating takes place at around 350 km where $Y_n \approx 0.53$, $Y_p \approx 0.47$ [10, 28], one arrives at the following estimate:

$$R_{RSFP} \approx 2.1. \tag{24}$$

It is interesting to compare this with the analogous simple estimate of the energy gain due to the MSW effect. For the same values of Y_p, Y_n and neutrino temperatures the gain factor would be $R_{MSW} \approx 2.5$[1], i.e. about 20% larger than in the case of the RSFP transition. There are two reasons for this: first, the MSW effect would convert ν_μ or ν_τ into ν_e increasing the mean energy of the electron neutrinos by the factor $T_{\nu_\mu}/T_{\nu_e} \approx 2$, whereas the RSFP transition $\nu_{\mu(\tau)L} \to \bar{\nu}_{eR}$ increases the energy of electron antineutrinos by a smaller factor, $T_{\nu_\mu}/T_{\bar{\nu}_e} \approx 1.7$; second, the ν_e's which play a major role in the MSW–enhanced shock reheating interact with more abundant neutrons whereas the $\bar{\nu}_e$'s playing a major role in the RSFP–enhanced reheating interact with less abundant protons. Comparison of the two mechanisms shows that in general the RSFP–enhanced shock reheating is likely to be more efficient for smaller values of Δm^2 since the adiabaticity parameter (9) becomes larger and the resonance occurs at lower densities, which means that the relative fraction of protons in the matter between the positions of the resonance and the stalled shock becomes

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[1]Fuller et al. [10] obtained $R_{MSW} \approx 2$ for slightly different values of neutrino temperatures, Y_p and Y_n.

larger. On the contrary, with increasing values of Δm^2 the resonance would occur at higher densities where the MSW transitions are more adiabatic and in addition the relative neutron abundance is higher, and so the MSW effect would become increasingly more important.

It should be emphasized that our conclusion that the RSFP–induced neutrino conversions are typically slightly less efficient in reheating the shock than the MSW effect was the direct consequence of our assumption that $\Delta m^2 > 0$. However, neutrino masses do not have to follow the pattern of the charged fermion mass hierarchy; it is quite possible that they have an inverse hierarchy in which the ν_e's are the heaviest among the neutrinos (see, e.g., [32, 33, 34] for particle–physics models and [35, 36, 33, 34, 37] for discussions of possible phenomenological consequences). The neutrino mass hierarchy can be even more complicated with, e.g., the ν_e mass being in between m_{ν_μ} and m_{ν_τ} (in this case the MSW effect or the RSFP can be responsible for the solar neutrino deficit, while for ν_e being the heaviest neutrino they would not be operative in the sun). In any case, if m_{ν_e} is larger than at least one of the other neutrino masses, spin–flavor conversion of the type $\bar{\nu}_\mu(\bar{\nu}_\tau) \to \nu_e$ will be resonantly enhanced in the supernova environment. This means that the resulting electron neutrinos will have a mean energy which is twice the energy of the originally produced ν_e's, and the gain in the shock reheating energy will be exactly the same as in the case of the MSW transitions. It should be noted, however, that because of the experimental upper limit $m_{\nu_e} < 4.35$ eV (95% c.l.) [38], the resonant $\bar{\nu}_\mu(\bar{\nu}_\tau) \to \nu_e$ conversion can only be relevant for supernova shock reheating for the electron neutrino mass lying in the relatively narrow range $\sim (3 - 4.3)$ eV. In addition, in this case the RSFP transition may be constrained severely by the supernova nucleosynthesis (r-process) arguments similar to those applied to the MSW effect in supernovae [28, 31, 39]. At the same time, the RSFP–induced neutrino conversion with direct neutrino mass hierarchy may even help the supernova nucleosynthesis process [40].

We would like to comment now on the implications of the RSFP–induced $\nu_\mu(\nu_\tau) \to \bar{\nu}_e$ conversion for the neutrino signals from supernovae. As we emphasized before, as a result

of this conversion the $\bar{\nu}_e$'s emerging from the supernova would have higher energies than the originally produced ones. This should result in a "stiffer" than expected spectrum of electron antineutrinos observed through the reaction $\bar{\nu}_e + p \rightarrow n + e^+$ in the terrestrial water Čerenkov detectors. Similar consequences would result for large-mixing-angle neutrino oscillations. Since the SN1987A $\bar{\nu}_e$ signals observed by Kamiokande and IMB detectors are in reasonable agreement with expectations, one might conclude that the conversions of $\nu_\mu(\nu_\tau)$ or their antiparticles into $\bar{\nu}_e$ are disfavored by the data [41]. However, this conclusion relies heavily on the theoretical predictions for the spectra of the supernova neutrinos and is therefore model dependent. In addition, though the SN1987A neutrino observations have confirmed the basic ideas of the supernova explosion and neutrino transport theory, the signals observed by the Kamiokande and IMB detectors are not fully understood. The opinions on whether a significant fraction of the SN1987A ν_μ's or ν_τ's or their antiparticles could have been converted into $\bar{\nu}_e$'s leading essentially to an interchange of their spectra differ significantly. The authors of [41] and [35] believe that such a possibility is essentially ruled out, whereas the authors of [36, 43, 42] conclude that there is no useful limit on the probability of such an interchange. The authors of ref. [44] turned the argument around and pointed out that if the solution of the solar neutrino problem through the MSW effect with the parameters that would lead to a significant interchange of the hard and soft neutrino spectra in the supernova is borne out by future solar neutrino experiments, one would have to conclude that the supernova $\nu_\mu(\nu_\tau)$ and/or $\bar{\nu}_e$ spectra are softer than had been thought previously.

Without entering this discussion, we would like to point out a mechanism which could lead to the $\bar{\nu}_e$'s observed in the terrestrial detectors having the expected or even a softer spectrum even if a strong RSFP–induced $\nu_\mu(\nu_\tau) \rightarrow \bar{\nu}_e$ transition occurs in the supernova. The possibility of having a softer than expected $\bar{\nu}_e$ spectrum is especially interesting since it is favored by the SN1987A data. It has been shown in [45] that, if the neutrinos have non-zero vacuum mixing angle in addition to the transition magnetic moment, and the direction of the transverse magnetic field changes along the neutrino path, resonant $\bar{\nu}_e \leftrightarrow \nu_e$

transitions are possible. In this case the $\bar\nu_e$'s observed in the terrestrial detectors may have the spectra of the originally produced ν_e's or $\bar\nu_e$'s, depending on the history of the conversions that the neutrinos experienced in the supernova before they reach the $\bar\nu_e \leftrightarrow \nu_e$ resonance. This possibility will be discussed in more detail elsewhere.

Our estimate of the energy gain due to the RSFP of supernova neutrinos was in fact very rough. More accurate calculations would involve integration of the energy deposition rate over the entire region between the resonance and the position of the shock, and should take into account neutrino re-emission by matter [the last term in eq. (22)]. Such a calculation was carried out in [10] for the MSW–enhanced shock reheating. The result was about a factor 0.8 decrease of the energy gain ratio R_{MSW} as compared with the simple estimate based on the formula similar to our eq. (23). The main reason for this decrease is that MSW–enhanced shock reheating increases the local temperature of the matter T_m and thus the emissivity term in eq. (22), resulting in a "negative feedback" for the effect [10]. We expect a similar reduction to take place for the RSFP–enhanced reheating, and therefore the estimate of eq. (23) should be replaced by $R_{RSFP} \simeq (1.6\text{–}1.7)$.

5   Conclusions

We have shown that resonant spin–flavor precession of neutrinos due to interaction of their transition magnetic moments with the strong magnetic fields inside supernovae may increase the energy deposited by neutrinos in the matter behind the shock by about 60% and thus help to re-accelerate it. For the process to be efficient in the supernova environment, the heavier neutrino mass should be in the range $\sim (3-600)$ eV, and the transition magnetic moment μ should be of the order of $10^{-14}\mu_B$ provided that the magnetic field strength at the resonance position is of the order of $10^{12}-10^{15}$ G. All these values of the parameters are consistent with the available laboratory and astrophysical constraints on neutrino properties as well as our present ideas about supernova magnetic fields. In fact, the neutrinos with the

masses and magnetic moments in the above range may be very interesting for cosmology and for the decaying neutrino theory [21].

Our simple consideration gave only a rough estimate of the supernova explosion energy increase due to the RSFP conversion; whether or not this effect is sufficient to revive the shock leading to a successful supernova explosion can only be decided on the basis of a full-scale supernova dynamics calculation with the RSFP transition included, which goes beyond the scope of the present paper. However, our estimates show that the RSFP–induced neutrino conversion can result in quite a sizable increase of the supernova explosion energy, and we believe that this effect deserves further investigation.

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Figure captions

Fig. 1. The characteristic scale height L_ρ as a function of the radial distance r from the center of the supernova at $t \approx 0.15$ s after the bounce. The shock wave is located at $r \approx 430$ km. The matter density and Y_e profiles of refs. [10, 28] have been used.

Fig. 2. The transition probability P versus $E/\Delta m^2$ for two different magnetic field distributions. The solid line corresponds to a power law [cf. eq. (13)] with $k = 2$ and $\mu B_0 = 5.8 \times 10^{-8}$, eV, the dashed line, to a power law with $k = 3$ and $\mu B_0 = 2.9 \times 10^{-7}$ eV.
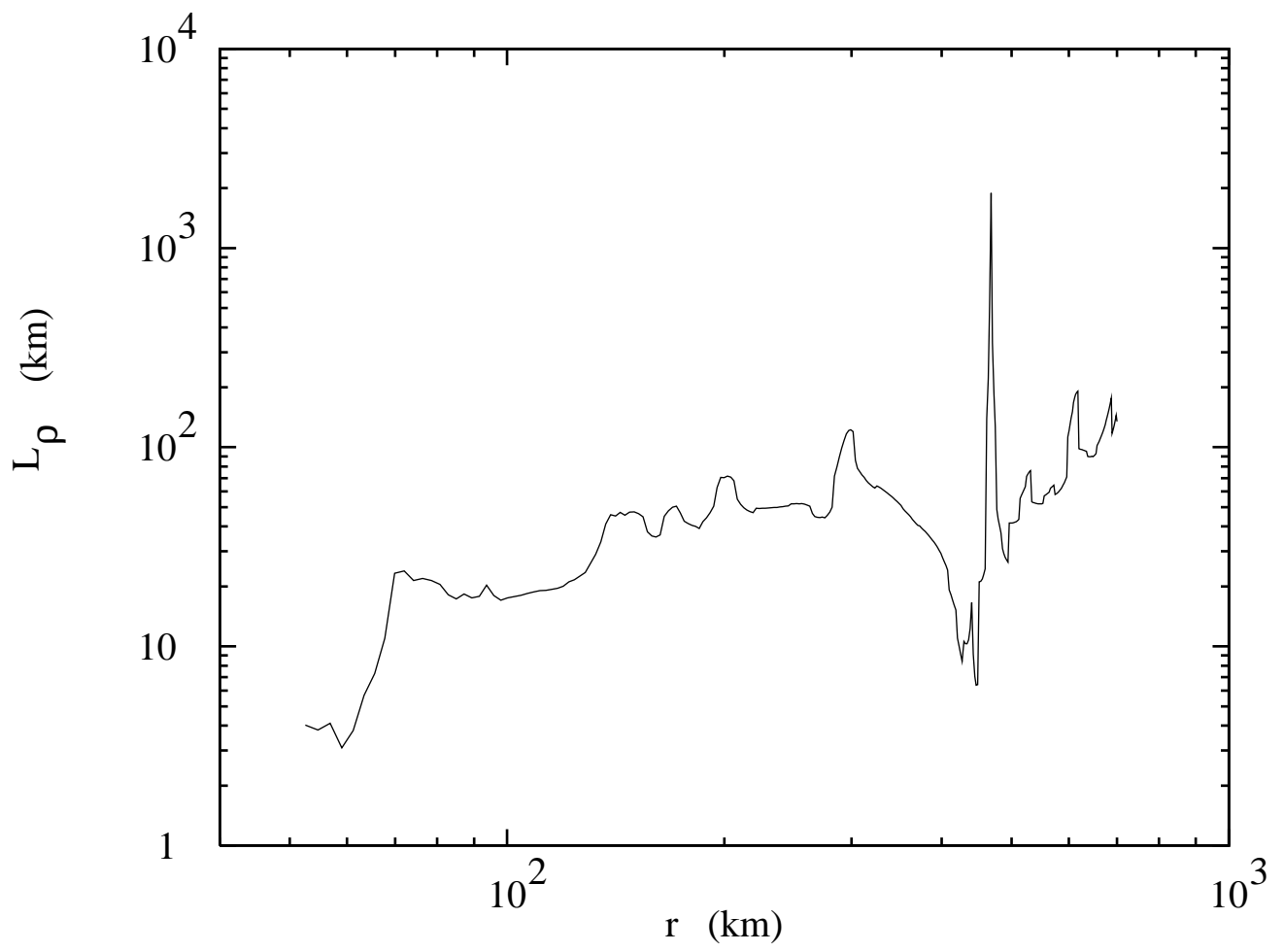
Fig. 1

Fig. 2